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SU(3) constraints on cryptoexotic pentaquarks

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We examine SU(3) constraints on the spectrum and decay widths of cryptoexotic nucleon-like states in antidecuplet of pentaquarks. We show that in the ideal mixing scenario the number of free matrix elements describing splittings reduces from 4 to 2. Model-independent sum rules are derived. Using as input Θ^+ and $\Xi_{3/2}$ masses we show that it is difficult to interpret Roper and $N^*(1710)$ nucleon resonances as cryptoexotic pentaquarks. Large N_c limit for antidecuplet and accompanying octet in the diqaurk picture and the analogy with the soliton model is also discussed.

1 Introduction

The discovery of the strange baryon Θ^+ at 1540 MeV [1] - [11] opened a new chapter in hadron spectroscopy. If confirmed by high statistics experiments, it will challenge the commonly used naive quark model that worked so well for almost four decades. Two questions will have to be answered: 1) what is the new dynamical mechanism overlooked so far that renders light and narrow exotic pentaquarks and 2) why the nonexotic states are not affected by this mechanism, in other words: why the naive quark model works so well in the nonexotic sector. Of course a likely answer might be that the key ingredient missing in the naive quark model is chiral symmetry and that baryons are solitons in an effective chiral model [12] - [14]. After all, exotic antidecuplet emerges naturally in chiral soliton models [15] - [18]. Indeed, light and narrow Θ^+ was predicted many years before its discovery [19] in the chiral quark-soliton model [20] - [23]. Even then some aspects of the soliton models have to be clarified and better understood [24] - [28].

It is the purpose of the present paper to examine if the SU(3) symmetry alone is able to explain, or – more modestly – to accommodate further exotics that follow inevitably from the discovery of Θ^+. Experimentally another exotic state, namely $\Xi^{--}_{3/2}(1860)$ (here subscript $3/2$ refers to isospin), has been already announced by NA49 experiment at CERN [29]. Theoretically, it is clear that there must be other cryptoexotic states which span antidecuplet of flavor SU(3). These states are cryptoexotic since unlike Θ^+ and $\Xi_{3/2}$ their quantum numbers can be constructed from 3 quarks only, however symmetry requires that in fact they contain an additional $q\overline{q}$ pair. These cryptoexotic states, and more precisely the cryptoexotic nucleon-like states, are the primary subject of the present note. In Section 2 we calculate mass spectrum in the moderately broken SU(3) symmetry and then postulate ideal mixing. Model-independent sum rules for the masses of exotic baryons are derived. We also show that the ideal mixing (*i.e.* requirement that physical states are defined as the ones which posses definite number of (anti)strange quarks) is – under assumption of small symmetry breaking – unable to accommodate known nucleon resonances such as Roper and $N^*(1710)$. In Sect. 3 we further show that also the decay widths are incompatible with ideal mixing.

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In Sect. 4 we slightly change the subject. In the large N_c limit, flavor SU(3) representations are no longer ordinary octet and decuplet. In soliton models their generalizations are straightforwardly constructed [28, 30, 31] because of the constraint coming from the Wess-Zumino term [32] - [34]: $8=(1,1)\rightarrow(1,(N_c-1)/2)$ and $10=(3,0)\rightarrow(3,(N_c-3)/2)$ [1]. Moreover $\overline{10}=(0,3)\rightarrow(0,(N_c+3)/2)$. We show that in the quark model with diquark correlations, the construction in which diquarks contain 2 quarks for arbitrary N_c reproduces the same sequence of SU(3) representations. As a result the ideal mixing can be generalized to arbitrary (odd) N_c.

We summarize in Sect. 5

Most of this work was done during my short visits at the Institut für Theoretische Physik II of the Ruhr-University in Bochum in January and June 2003. It is always a pleasure to visit this lively international group led by Klaus Goeke whose support and enthusiasm stay beyond the predictions which led to the discovery of pentaquarks. It is a great honor to dedicate to Klaus this set of remarks on the SU(3) nature of exotic states.

2 SU(3) symmetry and ideal mixing

Quantum numbers of Θ^+ require that its minimal quark content is $|uudd\overline{s}\rangle$. Two quarks can be either in flavor $\overline{3}$ or 6. Therefore possible representations for 4 quarks are contained in the direct products

$$\begin{aligned}
&\overline{3}\otimes\overline{3}=3+\overline{6},\\
&\overline{3}\otimes 6=3+15,\\
&6\otimes 6=\overline{6}+15'+15.
\end{aligned}\tag{1}$$

Here $15=(2,1)$ and $15'=(4,0)$. Adding an $\overline{3}$ antiquark yields

$$\begin{aligned}
3\otimes\overline{3}&=1+8,\\
\overline{6}\otimes\overline{3}&=8+\overline{10},\\
15\otimes\overline{3}&=8+10+27,\\
15'\otimes\overline{3}&=10+35,
\end{aligned}\tag{2}$$

Therefore $\left|q^4\overline{q}\right\rangle$ state can be in one of the following flavor representations:

$$\left|q^4\overline{q}\right\rangle\in 1,\,8,\,10,\,\overline{10},\,27,\,35.\tag{3}$$

Whether all representations (3) are allowed depends on the dynamics of a specific model.

Out of allowed representations (3) the lowest one including explicitly exotic states is $\overline{10}$ which appears in a direct product of 4 quarks in flavor $\overline{6}$ and an antiquark:

$$\bar{6}\otimes\bar{3}\rightarrow 8+\overline{10}\tag{4}$$

and is therefore inevitably accompanied by an octet (see Fig. 1). Since both representations may have spin $1/2$ it is reasonable to assume that in the first approximation the masses of 8 and $\overline{10}$ are degenerate. Unlike in the case of the ordinary octet and decuplet, pentaquark symmetry states (*i.e.* states which are pure octet or antidecuplet) do not have a unique quark structure. For example a proton-like state in antidecuplet and

[1] Here we denote SU(3) representations either by (p,q) where number of quark indices is p and antiquark indices q, or by dimension with "bar" if $q>p$.

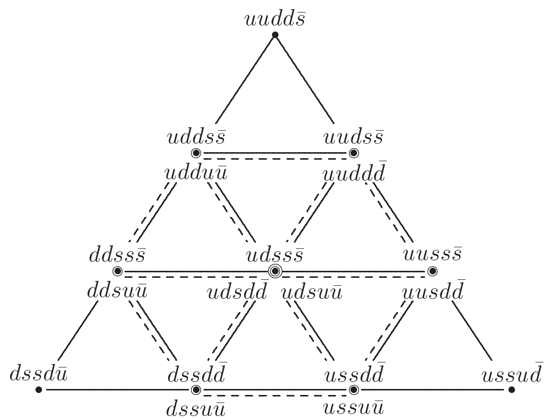

Fig. 1 Mixing of $\overline{10}$ and 8.

octet have the following quark content:

$$
\begin{aligned}
p_{\overline{10}} &= \sqrt{\frac{2}{3}}\,|uuds\overline{s}\rangle + \sqrt{\frac{1}{3}}\,\left|uudd\,\overline{d}\right\rangle , \\
p_{8} &= \sqrt{\frac{1}{3}}\,|uuds\overline{s}\rangle - \sqrt{\frac{2}{3}}\,\left|uudd\,\overline{d}\right\rangle
\end{aligned} \tag{5}
$$

where it is implicitly assumed that four quarks are in a pure $\overline{6}$ state. Similarly Σ-like states mix, while Θ^+ and Ξ-like states remain unmixed.

The QCD splitting hamiltonian ΔM that transforms like $Y=0$, $I=0$ component of the octet tensor operator, has the following matrix elements between the pentaquark states:

$$
\langle 8, YII_3\,|\Delta M|\,8, YII_3\rangle = D\left(I(I+1) - \frac{1}{4}Y^2 - 1\right) - FY,
$$

$$
\left\langle \overline{10}, YII_3\,|\Delta M|\,\overline{10}, YII_3\right\rangle = -CY \tag{6}
$$

and

$$
\left\langle 8, YII_3\,|\Delta M|\,\overline{10}, YII_3\right\rangle = \left\langle \overline{10}, YII_3\,|\Delta M|\,8, YII_3\right\rangle = \frac{G}{\sqrt{2}} \tag{7}
$$

for nucleon-like and Σ-like states. Constants D, F, C and G are unknown reduced matrix elements. Equations (6) would be the ordinary Gell-Mann–Okubo mass formulae if not for the mixing term (7) which is non-zero because pentaquark octet and antidecuplet have the same spin.

Now let us assume that strong interactions diagonalize (anti)strange quark content. The off-diagonal elements of the mass matrix take the following form for the nucleon-like states ($Y=1$):

$$
\langle udds\bar{s}|\,\Delta M\,|uddu\bar{u}\rangle = \langle uuds\bar{s}|\,\Delta M\,\left|uudd\bar{d}\right\rangle = \frac{\sqrt{2}}{6}(-G-2C+D+2F)
$$

and for $Y=0$, $I_3=\mp 1$ Σ-like states:

$$
\langle ddss\bar{s}|\,\Delta M\,|ddsu\bar{u}\rangle = \langle uuss\bar{s}|\,\Delta M\,\left|uusd\bar{d}\right\rangle = \frac{1}{3\sqrt{2}}(G-2D)
$$

and similarly for $I_3 = 0$. The requirement that mass matrix is diagonal in the (anti)strange quark basis leads to the relations:

$$G = 2D, \quad C = F - \frac{D}{2} \tag{8}$$

so that the whole spectrum depends on two constants D and F and the overall mass scale M:

$$\begin{aligned}
\Theta^+ &= M - 2F + D, \\
N^* &= M - F - \frac{3}{2}D, \\
N_s^* &= M - F + \frac{3}{2}D, \\
\Sigma^* &= M - D, \\
\Sigma_s^* &= M + 2D, \\
\Xi &= M + F - \frac{D}{2}.
\end{aligned} \tag{9}$$

Here by N^* and Σ^* we denote octet-like states with additional $u\overline{u}$ or $d\overline{d}$ pair and subscript s denotes additional $s\overline{s}$ pair. Ξ stands for both $I = 1/2$ and $I = 3/2$ states. The requirement that the mass matrix is diagonal in the (anti)strange quark basis reduced the number of the unknown constants describing the splittings from 4 to 2.

Scenario that the strong interactions diagonalize (anti)strange content is known as *ideal mixing*. It has been discussed in more detail in the context of the diquark model proposed by Nussinov [35] and Jaffe and Wilczek [36] where pentaquarks are viewed as bound states of two diquarks and an antiquark. Diqaurks are color and flavor antitriplets in a relative P-wave, hence two diquark state is a color 3 and flavor $\overline{6}$. The motivation for a diqaurk picture comes from the attraction between two quarks in a $\overline{3}$ color channel. Diquark correlations have been discussed in the context of meson spectroscopy [37]. A schematic hamiltonian for a diquark-like pentaquarks has been proposed in Ref. [36]. It contains two parameters: strange quark mass m_s and parameter α which is embodies the fact that diquarks involving strange quark are more tightly bound. These parameters are related to the reduced matrix elements introduced in (6) in the following way

$$F = \frac{1}{6}(5\alpha + 4m_s), \quad D = \frac{1}{3}(\alpha + 2m_s). \tag{10}$$

A somewhat different hamiltonian was introduced by Cohen in Ref. [38]. It is expressed in terms of a total number of strange quarks and the excess of strange over antistrange quarks. the hamiltonian $H = a(n_s + n_{\overline{s}}) + b' n_s$ also falls in the above category with $m_s \to a$ and $\alpha \to b'$ in (10).

A quantitative support for the dipole picture comes also from the instanton liquid model of the QCD vacuum. There, a tightly bound scalar diquark is found to have a mass of a single constituent quark, however, a tensor diquark in flavor 6 is only slightly heavier (570 MeV) [39]. If so, a bound state of scalar and tensor diquarks and a light antiquark would fall into $8 + 10 + 27$ of flavor SU(3) and, since there would be no penalty for the P-wave excitation, this state should be lighter than a bound state of two scalar diquarks in a P-wave and an antiquark. One should, however, consider this possibility with care, since once the diquarks overlap they may dissolve.

Any model hamiltonian of the SU(3) symmetry breaking with ideal mixing must obey equations (9). These equations lead to the model independent relations

$$\begin{aligned}
N_s^* - N^* &= \Sigma_s^* - \Sigma^*, \\
2(N_s^* - \Theta^+) &= \Xi - N^*, \\
N_s^* - \Sigma^* &= \Theta^+ - N^*
\end{aligned} \tag{11}$$

(and linear combinations of (11)).

Now we shall try to constrain further the remaining two free parameters F and D. To this end we plot in Fig. 2 the splittings (in units of D) as functions of parameter $x = F/D$. We see that for $1/2 < F/D < 5/2$ (depicted by thin vertical lines) the states are ordered according to the increasing number of (anti)strange quarks. For $F/D = 1$ the levels are equally spaced and the states with the same number of (anti)strange quarks are degenerate. It is therefore a reasonable starting point for a small perturbation in x which would lift this degeneracy introducing small splittings in a way similar to the $\Sigma - \Lambda$ splitting in the regular octet. We can see from Fig. 2 and Eqs.(9) that to the left of $x = 1$ the following splittings become equal: $\Theta^+ - \Sigma^* = \Xi - \Sigma^*$ for $x = 7/10$, whereas to the right of $x = 1$: $N_s - \Sigma^* = \Xi - N_s$ for $x = 3/2$. If the splittings of the states with identical (anti)strange content are to be smaller than those which differ by one (anti)strange quark, then one should require $7/10 < x < 3/2$. We shall shortly see that only $x > 1$ leads to the reasonable phenomenology.

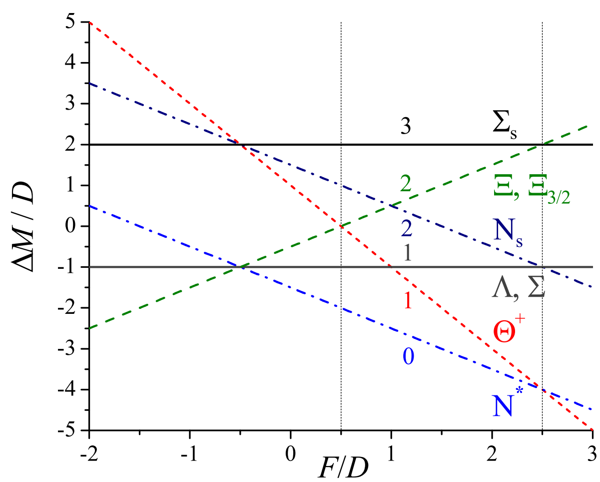

Fig. 2 Pentaquark spectrum in the ideal mixing scenario as a function of the parameter $x = F/D$. Numbers next to the lines denoting different particles refer to the total number of strange and antistrange quarks. For $1/2 < x < 5/2$ the states are ordered according to the number of (anti)strange quark content.

Now let us use experimental input for the masses of Θ^+ and $\Xi_{3/2}$. With this input the whole spectrum (9) depends only on one free parameter which we choose to be $x = F/D$. In Table (12) we display the masses of the remaining cryptoexotic states and the values of parameters D and F for three choices of x: 3/4, 1 and 3/2. We see that for $x \leq 1$ the lightest state, namely the N^* resonance, is unrealistically light. On the other hand for $x = 3/2$ the two nucleon-like resonances N^* and N_s^* have masses which allow to identify them with the Roper and $N^*(1710)$.

$x = F/D$	$3/4$	1	$3/2$
D [MeV]	427	213	107
F [MeV]	320	213	161
N^*	793	1222	1433
Σ, Λ	1327	1540	1647
N_s^*	2074	1860	1754
Σ_s	2608	2180	1968

(12)

Although this seems encouraging, two remarks are here in order. First, $x = 3/2$ is really the upper bound and the spectrum is equally spaced with $\Delta M \sim 107$ MeV. This contradicts our original assumption that perturbation around $x = 1$ should be small. The second remark concerns the decay widths and will be discussed in the next Section.

3 Decay widths in the ideal mixing scenario

The decay width of a symmetry state B in SU(3) flavor representation R into an octet baryon B' and a pseudoscalar octet meson φ reads:

$$\Gamma_{B\to B'\varphi} = \frac{G_R^2}{8\pi M\,M'} \times C^R(B,B',\varphi) \times p_\varphi^3. \tag{13}$$

Here M, M' and p_φ are baryon masses and meson momentum in the decaying baryon reference frame. $C^R(B,B',\varphi)$ is the relevant SU(3) Clebsch-Gordan coefficient (squared) and G_R the decay constant. For states which are mixtures of different representation states, like our candidates for Roper and $N^*(1710)$:

$$\begin{aligned} \text{Roper} \to |N_1^*\rangle &= \sqrt{\frac{1}{3}}\left|\overline{10},N\right\rangle - \sqrt{\frac{2}{3}}\left|8,N\right\rangle, \\ N^*(1710) \to |N_2^*\rangle &= \sqrt{\frac{2}{3}}\left|\overline{10},N\right\rangle + \sqrt{\frac{1}{3}}\left|8,N\right\rangle, \end{aligned} \tag{14}$$

the amplitude for the decay width would be the sum over $R = \overline{10}$ and 8. However, since the decay width of $\Theta^+ = \left|\overline{10},\Theta^+\right\rangle$ is very small, indicating that $G_{\overline{10}} \sim 0$ (remember Θ^+ does not mix), the $\overline{10}$ component in Eqs.(14) can be safely neglected in the first approximation. Then

$$\frac{\Gamma_{N_1^*\to N\pi}}{\Gamma_{N_2^*\to N\pi}} \sim 2\frac{M_2}{M_1}\frac{p_1^3}{p_2^3} \sim 0.75 \tag{15}$$

where we have used physical masses for Roper and $N^*(1710)$. Equation (15) indicates that partial decay widths of the nucleon-like pentaquark resonances have to be of the same order: either both small or both large. This observation was first made by Cohen [38] and presented in form of the inequality connecting different decay constants. Experimentally [38] $\Gamma_{\text{Roper}\to N\pi} \sim 228$ MeV and $\Gamma_{1710\to N\pi} \sim 15$ MeV. Order of magnitude difference between the partial decay widths of Roper and $N^*(1710)$ remains in contradiction with the ideal mixing scenario.

Of course ideal is mixing probably an idealization. The nonideal mixing would invalidate Cohen's inequality. More general scenarios have been discussed in Refs. [40] - [44]. However, the discussion of masses and decay widths of the N^* states under assumption that they correspond to the Roper and $N^*(1710)$ done in Ref. [41] still indicates that it is impossible to match the mass splittings with the observed branching ratios for these two resonances even for arbitrary mixing. It is shown that the *nonideal* mixing required for the decay $N^*(1710) \to \Delta\pi$ is not compatible with the mixing deduced from the masses. The conclusion of Ref. [41] is that most probably $\Xi_{3/2}(1860)$ is not a member of $\overline{10}$ but rather of 27. That would release constraints on the $N_{\overline{10}}$ state coming from the equal spacing of antidecuplet. Another possibility based on the nonideal mixing scenario advocated by Diakonov and Petrov [40] is that there should be a new N^* resonance in the mass range of $1650 \div 1680$ MeV, a possibility discarded in Ref. [41].

4 Large N_c limit: solitons vs. diquarks

For arbitrary N_c ordinary baryons consist of N_c quarks and "pentaquarks" are built from $N_c + 1$ quarks and an antiquark. In the soliton picture relevant flavor representations are selected by the requirement that

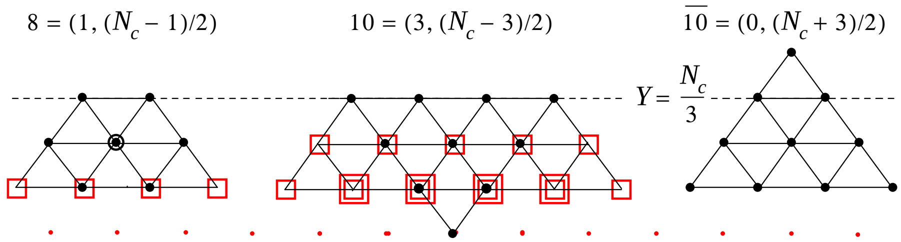

Fig. 3 SU(3)-flavor representations for arbitrary N_c. Representations are selected by a demand that they contain states with $Y = N_c/3$. Spurious states which disappear for $N_c = 3$ are denoted by open squares.

physical multiplets contain states with hypercharge $Y = N_c/3$, leading to the following generalizations [28, 30, 31]:

$$\begin{aligned} 8 &= (1,1) \to (1, \frac{N_c-1}{2}), \\ 10 &= (3,0) \to (3, \frac{N_c-3}{2}), \\ \overline{10} &= (0,3) \to (0, \frac{N_c+3}{2}) \end{aligned} \tag{16}$$

as illustrated in Fig. 3.

Although it possible to show on general grounds that representation content of the quark model and soliton model coincide for large N_c [16, 45], it is instructive to illustrate this in the specific quark model. Moreover, the general arguments apply to the highest representations, *i.e.* antidecuplet in the case of "pentaquarks" but not to the accompanying octet. In the following we shall construct large N_c generalizations of the flavor representations for "pentaquarks" in the diquark picture and show that they coincide with (16).

For "pentaquarks" in the diquark picture we would have $(N_c+1)/2$ diquarks (remember N_c is odd). For $N_c > 3$ two quarks antisymmetrized in color and flavor still form flavor $\overline{3}$ but color representation is (anti) $N_c(N_c-1)/2$. Here we denote SU(3) representations by the number of quark and antiquark indices (p,q) and/or for SU(N_c) by dimension with a suffix anti- (or bar) if $q > p$. Generalizing scenario of Jaffe and Wilczek [36] we put all diquarks in a color and space antisymmetric state. We can easily antisymmetrize $N_c - 1$ quarks leaving the last diquark aside. Adding this last diquark (antisymmetrized quark pair) results in two color structures, one of them being N_c. This is depicted in Fig.4. So adding an antiquark in color $\overline{N}_c$ will result in a color singlet state.

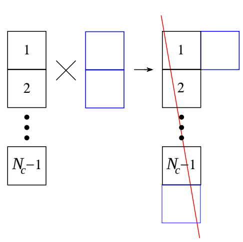

Fig. 4 Color representation for $(N_c+1)/2$ diquarks. First, $N_c - 1$ quarks are symmetrized and then the last two quarks are added in a form of a diquark to produce fundamental representation of SU(N_c).

In flavor space we have to symmetrize all diquarks forming representation $(0,(N_c+1)/2)$. Adding an antiquark in flavor $\overline{3} = (0,1)$, as depicted in Fig. 5, results in two representations: "8" $= (1,(N_c-1)/2)$ and "$\overline{10}$" $= (0,(N_c+3)/2)$. These are the same representations as in the soliton case (16).

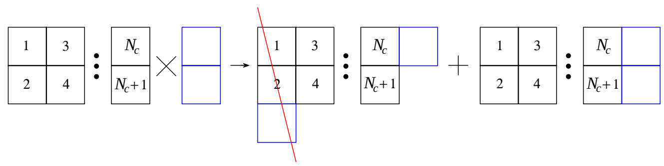

Fig. 5 Flavor representations for $(N_c+1)/2$ diquarks and an antiquark. Diquarks are symmetrized and adding an antiqauark produces "8" $=(1,(N_c-1)/2)$ and "$\overline{10}$" $=(0,(N_c+3)/2)$.

Note that this generalization of octet and antidecuplet allows for ideal mixing scenario. Indeed, both nucleon-like and Σ-like states mix in the same way as for $N_c=3$. Only $\Xi_{3/2}$ in "$\overline{10}$" will additionally mix with spurious sates in "8", however, the latter states will disappear for $N_c=3$.

Finally let us remark, that another generalization of the diquark picture which parallels junction models would result in a completely different representation structure. Indeed, another possible generalization of a diquark would be to antisymmetrize N_c-1 quarks. We would have then to take N_c-1 of such "diquarks" and add an antiquark. In this case "pentaquarks" would consist of $(N_c-1)^2$ quarks and an antiquark. This would be a completely different object than the regular pentaquark constructed from N_c+1 quarks and an antiquark, although accidentally for $N_c=3$ both pictures coincide.

5 Summary

Present study shows that consistent description of cryptoexotic nucleon-like states requires new N^* resonances in the mass range of $1650 \div 1680$ MeV. Similar conclusion has been reached in Ref. [42] where mixing scenario was discussed within the framework of the quark soliton model. Here already the ordinary nucleon state has a non-negligible admixture of $\overline{10}$. Taking this mixing into account the authors of Ref. [42] estimated the width of a possible but yet undiscovered N^* state of a mass of 1680 (1730) MeV to be $\Gamma_{N^*\to N\pi} \sim 2.1$ (2.3) MeV. Although quite narrow, this width is too large to be accommodated in the πN scattering data. The same authors [42] claim that the improved phase shift analysis admits two candidates for the narrow resonances of these masses but with decays widths smaller than 0.5 (0.3) MeV. Further decrease of theoretical predictions might be achieved by adding a mixing to yet another nucleon-like state as Roper and $N^*(1710)$. And even further suppression of this decay is provided by the 27 admixture as discussed in Ref. [43, 44].

We have also shown that a particular construction of SU(3) flavor baryon representations that generalizes the diquark model for large N_c, coincides with the soliton model.

To conclude let us note that physics of N^* and Σ^* states will be most probably physics of extensive mixing between different nearby states. It is clear from our discussion above that the consistent physical picture requires the existence of yet undiscovered nucleon resonances within the mass range of $1650-1750$ MeV or so. Experimental searches for new narrow nucleon-like states have been recently performed by some experimental groups with positive preliminary evidence [46, 47]

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